Strange Dibaryon Production in Relativistic Heavy Ion Collisions

C. B. Dover

Physics Department
Brookhaven National Laboratory
Upton, New York

BNL--41934

DE89 003159

ABSTRACT

We review theoretical predictions for the mass spectrum of strange dibaryons, with emphasis on $S = -2$ states. In a version of the Callen-Klebanov model applied to dibaryons, the lowest-lying configuration is an $SU(3)$ flavor $\{10^*\}$ with $J^P = 0^+$ and isospin $I = 1$, in contrast to the six-quark (Q^6) bag model, which yields the H dibaryon (flavor $\{1\}$, $J^P = 0^+$, $I = 0$) as the lowest $S = -2$ excitation. Possible experimental consequences of this different level order are discussed. For the H , we discuss formation rates in relativistic heavy ion collisions. For Brookhaven AGS energies (15 GeV/A), we find a rate of order 10^{-2} H 's per central collision. Since the H is produced in a high multiplicity environment in heavy ion collisions, its detection poses a problem. Several possible detection schemes for the H are investigated. These include studies of weak decays ($\Sigma^- p$ "vees"), diffractive dissociation ($H p \rightarrow \Lambda \Lambda p \rightarrow 3p + 2\pi^-$) and nuclear fragments with anomalous charge/mass ratios.

Invited Talk at
1988 International symposium on
Hypernuclear and Low-Energy Kaon Physics
Legnaro (Padua), Italy
September 12 - 16, 1988

Proceedings to be published in Nuovo Cimento

This manuscript has been authored under contract number DE-AC02-76CH00016 with the U.S. Department of Energy. Accordingly, the U.S. Government retains a non-exclusive, royalty-free license to publish or reproduce the published form of this contribution, or allow others to do so, for U.S. Government purposes.

MASTER

DISTRIBUTION OF THIS DOCUMENT IS UNLIMITED

ps

PACS. 14.20.Pt – Dibaryons

PACS. 25.70.-z – Heavy-ion-induced reactions

PACS. 25.70.Np – Fragmentation and relativistic collisions

1. – Introduction

The spectroscopy of multi-quark systems $Q^m \bar{Q}^n$ with $m+n \geq 4$ has provoked considerable interest, on both the theoretical and experimental sides. There is as yet no conclusive evidence for long-lived objects of such structure, although, as we argue below, the key experiments (i.e., for strangeness $S = -2$ dibaryons) have not yet been performed. The theoretical issue is the ability of dynamical models (quark bag, soliton, etc.) to predict the masses and decay modes of multi-quark states beyond the familiar baryons (Q^3) or mesons ($Q\bar{Q}$).

Enormous effort has been focused on the search for non-strange dibaryon resonances, since these are accessible through the nucleon-nucleon (NN) or $N\Delta$ channels. We do not review this work here, but rather devote our attention to $S = -2$ dibaryons. For certain quantum numbers, the addition of one or more strange (s) quarks to an assembly of up (u) and down (d) quarks is particularly favorable energetically, and thus the existence of stable or at least long-lived dibaryons becomes a distinct possibility. Some existing proposals to find $S = -2$ dibaryons employ the double strangeness exchange (K^-, K^+) reaction. The rates are low because of the limited intensity of available (or soon to be constructed) K^- beams. Here, we propose an alternative strategy, which utilizes the observed copious production of strange particles in relativistic heavy ion central collisions to produce $S = -2$ dibaryons via a fusion/coalescence mechanism.

In Chap. 2, we review theoretical attempts to predict the mass spectrum of $S = -2$ dibaryons, including the six-quark (Q^6) bag model and the soliton model à la Skyrme. These treatments produce quite distinct predictions for the level order, and different consequences for the dibaryon decay modes.

In Chap. 3, we specialize the discussion to the $0^+, I = 0$ H dibaryon, and we present some estimates of the production rate of the H from an hadronic fireball. For typical heavy ion collisions at the Brookhaven AGS, 15 GeV/A, the rate exceeds 1% per central collision. Because of the high multiplicity of particles (principally nucleons and pions), it is necessary to detect the H directly (unlike the (K^-, K^+) case, where one has two- or three-body kinematics, and it is sufficient to detect particles recoiling against the H). In Chap. 4, we investigate several possibilities for detecting the H. These include a search for the weak decay mode $H \rightarrow \Sigma^- p$, the dissociation of the H via the process $H + p \rightarrow \Lambda + \Lambda + p$, and the detection of nuclear fragments of anomalously low charge/mass ratio, which could arise as bound states of the H and a light nuclear core. Some conclusions are presented in Chap. 5.

2. – Spectroscopy and Decays of $S = -2$ Dibaryons

A variety of predictions for the mass spectra of strange dibaryons are available. Most early calculations are based on some version of the quark bag model⁽¹⁻⁶⁾ or the Skyrme-type soliton model⁽⁷⁻¹²⁾. A review of this work can be found in ref. (13). More recently, the H dibaryon mass has been calculated in lattice QCD^(14,15), as well as in the quark cluster/resonating group method^(16,17). Weakly bound deuteron-like $S = -2$ dibaryons can also exist in meson exchange models⁽¹⁸⁾. In the most complete lattice QCD⁽¹⁵⁾ and quark cluster⁽¹⁷⁾ calculations, the H dibaryon⁽¹⁾ lies below the $\Lambda\Lambda$ threshold, and is thus stable against strong decay. However, the predicted binding energy $B_H \equiv 2m_\Lambda - m_H$ is strongly model dependent: the converged lattice QCD result⁽¹⁵⁾ gives a deeply bound H ($m_H \approx 2m_n$), while the quark/gluon plus meson exchange model of Straub *et al*⁽¹⁷⁾ yields a weakly bound system ($B_H \approx 20$ MeV).

In addition to the H, the $S = -2$ dibaryon spectrum contains a number of excited states. These have been worked out in the context of the MIT Bag Model, for instance by the Nijmegen group⁽²⁾. The spectrum from Mulders *et al*⁽²⁾ is displayed in Fig. 1. Except for the H, the various configurations lie above one or more strong decay thresholds. In the bag model for $S = -2$, the masses $M(J, \{f\})$ for spin J and flavor representation $\{f\}$ are given by

$$M(J, \{f\}) = m_H + \Delta M \left[\frac{J(J+1)}{3} + f^2 \right] \quad (1)$$

where $m_H = 2160$ MeV/c², $\Delta M = 50$ MeV (ref. (2)) and f^2 is the eigenvalue of the usual quadratic Casimir operator F^2 of SU(3): we have $f^2 = (0, 3, 6, 8, 12, 18)$ for the $\{f\} = \{1, 8, 10 \text{ or } 10^*, 27, 35, 28\}$ dimensional representations of SU(3), respectively. The H is the $J = 0$, $I = 0$ SU(3)-flavor singlet, for which the color magnetic attraction due to quark/gluon exchange is maximized (i.e., f^2 is minimum for $\{f\} = \{1\}$). One should keep in mind the possibility that, if the H is more deeply bound than in Fig. 1, some of the excited states could also be stable (or at least long-lived). In Table 1, we give the s -wave ($L = 0$) baryon-baryon (BB) wave function components for the states in Fig. 1, along with the BB decay modes which are energetically allowed. For $H''_{0,1}$, the decay modes $\Xi N \pi$ or $\Lambda \Lambda \pi$ are also allowed, but the available phase space is small (if one believes the masses in Fig. 1!). Since $H''_{0,1} \rightarrow BB\pi$ requires a p -wave pion relative to the BB pair, the decay rate should be suppressed. Note that for H'_1 the $\pi^\pm H$ (1P_1) decay mode is allowed (but presumably suppressed because of the relative p -wave), while H'_0 , $H''_{0,1}$ cannot decay to πH because of spin-parity conservation. In Table 1, $\Sigma^* = \Sigma(1385)$, $\Xi^* = \Xi(1530)$, with decays $\Sigma^* \rightarrow \Lambda \pi$ (88%), $\Sigma \pi$ (12%) and $\Xi^* \rightarrow \pi \Xi$ (100%).

Consider now the states $H''_{0,1}$ (the subscript refers to isospin), which have intrinsic spin two, and decay via spin flip transitions into $L = 2$ baryon-baryon channels. The resultant $^5S_2 \rightarrow ^1D_2$, 3D_2 transitions have $\Delta L = 2$, $\Delta S = 1, 2$ and hence can only be mediated by the tensor operator S_{12} , generated by π or $\bar{K}$ exchange, for instance via the sequence $H''_0 \rightarrow \Sigma^* \Sigma (L = 0) \xrightarrow{\pi} \Lambda \Lambda (L = 2)$ or $H''_0 \rightarrow \Xi^* N (L = 0) \xrightarrow{\bar{K}} \Lambda \Lambda (L = 2)$. Using the masses

from Fig. 1, the c.m. momentum release k in the decays $H_{0,1}'' \rightarrow BB$ is given by 420 MeV/c for $H_{0,1}'' \rightarrow \Xi N$, 455 MeV/c for $H_0'' \rightarrow \Lambda\Lambda$, 340 MeV/c for $H_1'' \rightarrow \Lambda\Sigma^\pm$ and 175 MeV/c for $H_0'' \rightarrow \Sigma\Sigma$. The first three of these transitions are well matched kinematically to a $\Delta L = 2$ operator, i.e. $\Delta L \approx kR$, where $R \approx 1$ fm. Thus, such spin-flip decays are not (*a priori*) suppressed, as suggested in ref. (2). It would be useful to perform quantitative estimates of the decay widths for the excited $H_{0,1}'$ and $H_{0,1}''$ states. If they are sufficiently narrow, they could show up in plots of invariant mass for the $\Lambda\Lambda$, $\Lambda\Sigma$, $\Sigma\Sigma$ or ΞN systems.

The Skyrme model (19) was recently revived by Witten (20), who provided a strong motivation for it in terms of the large N_c limit of QCD. The model was extended to $B = 2$ systems in refs. (7-9); here all $SU(3)$ coordinates were treated as collective modes, so that the pion and kaon fields were treated on the same footing. This extension of the $SU(2)$ Skyrmion to $SU(3)$ fails to reproduce the $S = -1$, $B = 1$ spectrum, unless one adopts an unacceptably small value of the pion decay constant F_π . A key advance was made by Callen and Klebanov (11), who incorporate strong $SU(3)$ breaking from the beginning by treating the K and π fields on a separate footing: the Lagrangian is expanded to second order in the K field (vibrational approximation), while the π field retains a classical limit. The $B = 1$, $S = -1$ configurations then correspond to kaon bound states in the background field provided by the $SU(2)$ Skyrmion (11).

For $B = 1$, the Wess-Zumino term (11)

$$\mathcal{L}_{WZ} = \frac{2iN_c}{5F_\pi^2} B^\mu \left[(D_\mu K)^+ K - K^+ D_\mu K \right], \quad (2)$$

where B^μ is the $SU(2)$ baryon current and K is the kaon field, plays a crucial role. It yields the observed baryon multiplet structure ($\{8\}$ and $\{10\}$), through a constraint $Y_R = N_c n/3$ on the body-fixed hypercharge Y_R (n = number of quarks). In addition, $\mathcal{L}_{WZ}$ serves to split the $S = \pm 1$ mass eigenvalues. The "exotic" $S = +1$ baryons are pushed to high mass, a very desirable phenomenological feature.

The Skyrme model à la Callen-Klebanov has been extended to $B = 2$ systems by Kunz and Mulders (12). Unlike the earlier $SU(3)$ generalization of refs. (7-9), which yielded a $J^\pi = 0^+$, $I = 0$ configuration as the $S = -2$, $B = 2$ ground state (like the six-quark bag model), the treatment of the K field as a vibration leads to a lowest-lying state $\tilde{H}$ with $J^\pi = 0^+$, $I = 1$, $\{f\} = \{10^*\}$ and mass $m_{\tilde{H}} \approx 2350$ MeV/c². The first excited state $\tilde{H}'$ has $J^\pi = 0^+$, $I = 0$ (which corresponds to the quantum numbers of the H in the Q^6 model), but $\{f\} = \{27\}$ rather than $\{f\} = 1$. In fact, in the Skyrme model, the constraint on Y_R restricts the $B = 2$, $S = -2$ system to flavor $\{f\} = \{10^*\}$, $\{27\}$, $\{35\}$. The low-dimensional representations $\{f\} = \{1\}$ and $\{8\}$, which correspond to the states H , $H_{0,1}'$, $H_{0,1}''$ in the Q^6 spectrum (2) of Fig. 1, are absent in the Skyrme model spectrum (12). The mass spectrum of Kunz and Mulders (12) may be written in the form

$$M(J, I, \{f\}) = M_0 + M_1 J(J+1) + M_2 I(I+1) + M_3 \kappa(\kappa+1) + M_4 (1 - (-)^K) \quad (3)$$

where the parity is $(-)^K$ and the constants are given by $M_0 \approx 2300$, $M_1 \approx 38$, $M_2 \approx 24.5$, $M_3 \approx 34$, and $M_4 \approx -68$ (all in units of MeV/c^2). In comparison to Eq. (1), the above mass formula contains an explicit term $M_2 I(I+1)$ which breaks isospin degeneracy, and a “rotational” term in the flavor index κ (0, 1, 2 for $\{10^*\}$, $\{27\}$, $\{35\}$, respectively). Note also that for fixed (J, I, κ) , the negative parity states ($K = \pm 1$) lie *below* those of positive parity (ex., 2^- below 2^+).

In contrast to the Q^6 spectrum of Fig. 1, all masses in Eq. (3) lie above the $\Lambda\Lambda$, ΞN and $\Lambda\Sigma$ thresholds. Thus, one might think that all the states in the $B = 2$, $S = -2$ Skyrme spectrum ⁽¹²⁾ are broad. This is incorrect, since the generalized Pauli principle (GPP) acts to strongly restrict the allowed decays. The GPP implies the following selection rules ⁽²⁾:

$$\begin{aligned} 0^+, 2^+ \{8_A, 10, 10^*\} &\not\rightarrow B_8 B_8 \\ 1^+ \{1, 8_S, 27\} &\not\rightarrow B_8 B_8 \\ 0^+, 2^+ \{10^*, 35\} &\not\rightarrow B_{10} B_{10} \\ 1^+ \{27, 28\} &\not\rightarrow B_{10} B_{10} \end{aligned} \quad (4)$$

where B_8 and B_{10} are any members of the $SU(3)$ baryon octet or decuplet, respectively. Applying Eq. (4) to the ground state $\tilde{H}(0^+, \{10^*\})$ of the Skyrme spectrum ⁽¹²⁾, we obtain

$$\tilde{H} \not\rightarrow \Lambda\Sigma, \Xi N, \Sigma^* \Sigma^*, \Delta \Xi^* \quad (5)$$

In addition, the decays $\tilde{H} \not\rightarrow B_8 B_{10}$ ($L = 0$) are forbidden by spin-parity conservation. Thus, $\tilde{H}$ can only enjoy virtual dissociation through $L \neq 0$ partial waves, for instance

$$\tilde{H} \rightarrow \begin{cases} [\Sigma(1385) \otimes \Lambda]_{35} D_0 \\ [\Sigma(1750) \otimes \Lambda]_{33} F_0 \end{cases} \quad (6)$$

where $\Sigma(1750)$ is a $1/2^-$ baryon decaying principally to $\bar{K}N$ (10 – 40%) and $\Sigma\eta$ (15 – 55%). However, if $m_{\tilde{H}} \approx 2350 \text{ MeV}/c^2$, as per ref. ⁽¹²⁾, then both the $\Lambda\Lambda\pi$ ($2370 \text{ MeV}/c^2$) and $\bar{K}N\Lambda$ ($2550 \text{ MeV}/c^2$) decay channels are closed, and the transitions (6) cannot go on-shell. The conclusion is that $\tilde{H}$ would be stable against strong decays if $m_{\tilde{H}} < 2370 \text{ MeV}/c^2$.

The quark model and Skyrme spectra, which have the same quantum number content for $B = 1$, differ radically for $B = 2$, as illustrated in Table 2. The Q^6 spectrum for $S = -2$ contains no “extraneous” states ⁽²¹⁾, i.e., those which have no coupling to $B_8 B_8$ channels, while the Skyrme spectrum ⁽¹²⁾ exhibits several of these, including the 0^+ , $I = 1$, $\{10^*\}$ ground state $\tilde{H}$.

3. – Production of the H Dibaryon in Heavy Ion Collisions

Numerous methods have been suggested for producing the H (for a review, see ref. ⁽¹³⁾). There are two approved experiments for the Brookhaven AGS ⁽²²⁾ and another under way

at KEK in Japan (²³) which utilize the double strangeness exchange (K^- , K^+) reaction to produce an $S = -2$ system. In AGS 813, the $K^- + p \rightarrow K^+ + \Xi^-$ reaction is used to tag the formation of a Ξ^- , which is then slowed down electromagnetically and captured in deuterium. One then looks for a mono-energetic neutron from the process $\Xi^- + d \rightarrow n + H$ (²⁴). In AGS 835, the reaction $K^- + {}^3\text{He} \rightarrow K^+ + H + n$ is studied (²⁵). A status report on these two (K^- , K^+) experiments has been given at this conference by P. Barnes (²⁶). It has also been proposed (²⁷) that the H may be found in $\bar{p}$ -nuclear reactions. The case of $\bar{p} + {}^3\text{He} \rightarrow (KK\pi)^+ + H$ was discussed in detail by Guaraldo (²⁸) at this conference. Note also that an event which could be interpreted as the formation and decay of an H has been reported by Shahbazian et al (²⁹).

It has recently been proposed by Dover, Koch and May (³⁰) that relativistic heavy ion collisions would be a copious source of H dibaryons. Estimates of the production probability for the H have been given for the case of a hot baryon-rich hadron gas and also for a quark gluon plasma (³⁰).

For Brookhaven AGS energies (15 GeV/nucleon, fixed target), calculations in the hadron basis are appropriate. We assume that the complicated multi-body interactions in the HI collision result in the formation of a hot, dense hadron fireball, nearly at rest in the c.m. system of the colliding nuclei. During this stage, a substantial number of kaons (K) and hyperons (Y) are produced via the $NN \rightarrow NKY$ and $\pi N \rightarrow KY$ associated production reactions. For a ${}^{28}\text{Si} + {}^{197}\text{Au}$ collision at 15 GeV/A, preliminary experimental results (³¹) indicate that about 5 K^+ 's are formed in a central collision.

We can understand this result as follows: for a projectile (A_p) - target (A_T) collision in the stopping regime, the number of participants N is given geometrically as

$$N \approx A_p + \frac{3}{2} A_p^{2/3} A_T^{1/3} \approx 100 \quad (7)$$

for ${}^{28}\text{Si} + {}^{197}\text{Au}$, where the second term represents the number of target nucleons in a tube cut out by the projectile. The number N_p of participant protons is then $N_p = N/2 \approx 50$. From the observed (³¹) ratios $N_p/N_{\pi^+} \approx 5/2$ and $N_{K^+}/N_{\pi^+} \approx 1/4$, we then obtain $N_{K^+} \approx 5$. Further, using $N_{K^-} \approx 1$ and assuming $N_{K^0} = N_{K^+}$, $N_{\bar{K}^0} = N_{K^-}$, $N_{\Sigma^+} = N_{\Sigma^-} = N_{\Lambda}$, we can invoke strangeness conservation (in the absence of $\bar{\Lambda}$'s) to write

$$N_{\Lambda} \approx \frac{1}{2} (N_{K^+} - N_{K^-}) \approx 2 \quad (8)$$

Thus, in a typical ${}^{28}\text{Si} + {}^{197}\text{Au}$ central collision, one $\Lambda\Lambda$ pair is produced.

The H can now be produced from BB fusion in the hot hadron gas, according to the wave function

$$\Psi_H = \frac{1}{\sqrt{8}} (\Lambda\Lambda + \Sigma^0\Sigma^0 + \Sigma^+\Sigma^- + \Sigma^-\Sigma^+ + \Xi^0n + n\Xi^0 - \Xi^-p - p\Xi^-) \quad (9)$$

We use the relative phases from refs. (^{13,25}); this choice does not influence our results, since only probabilities for various BB channels enter the fusion calculation. The issue of

overall normalization is more critical: Donoghue *et al* ⁽³²⁾ as well as Gonzalez and Vento ⁽³³⁾ have pointed out that a complete basis for expanding the H (or, in general, any color singlet state) wave function is provided by a direct product of color singlet baryon states. The $8_c \otimes 8_c$ hidden color states (which usually are taken to constitute $\sqrt{4/5}$ of the H wave function, in amplitude) can be written in terms of products of color-singlet hadronic quark-cluster states, and hence they constitute another representation of the same wave function, rather than a component which is necessary to form a complete basis. Thus $\sqrt{1/8}$ rather than $\sqrt{1/40}$ appears as the normalization factor in Eq. (9).

The formation of the H is treated in terms of the fusion of two Q^3 bags into a Q^6 bag. Following ref. ⁽²⁵⁾, the fusion amplitude A_{ij} for a baryon pair $B_i B_j$ is given by the momentum space overlap

$$A_{ij} = \frac{1}{(2\pi)^{18}} \int \frac{d^3 \underline{k}_1}{2E_1} \cdots \frac{d^3 \underline{k}_6}{2E_6} \psi_H(\underline{k}_1 \cdots \underline{k}_6) \psi_i(\underline{k}_1, \underline{k}_2, \underline{k}_3) \psi_j(\underline{k}_4, \underline{k}_5, \underline{k}_6) \quad (10)$$

For Gaussian wave functions, one obtains

$$A_{ij} = (2\pi)^3 (2E_H 2E_i 2E_j)^{1/2} \delta^{(3)}(\underline{k}_H - \underline{k}_i - \underline{k}_j) \Gamma_{ij}(\underline{k}_i - \underline{k}_j) \quad (11)$$

$$\Gamma_{ij}(\underline{k}) = C_{ij} \left(\frac{8\pi R_H^2}{3} \right)^{3/4} e^{-R_H^2 k^2/12}$$

where C_{ij} is the color-spin-flavor factor ($1/8$ for $\Lambda\Lambda$, $3/8$ for $\Sigma\Sigma$, $1/2$ for ΞN), and we have used the approximation $R_H = R_i = R_j$ for simplicity ($R_H = \text{r.m.s. radius of the H}$).

We now obtain the total number N_H of H's (in the rest system of the fireball) by overlapping the fusion probability $|A_{ij}|^2$ with Boltzmann momentum distributions f_i and f_j for the baryons:

$$N_H = \sum_{i,j} \int \frac{d^3 \underline{p}_i}{(2\pi)^3} \int \frac{d^3 \underline{p}_j}{(2\pi)^3} \int \frac{d^3 \underline{p}_H}{(2\pi)^3} f_i(\underline{p}_i) f_j(\underline{p}_j) |A_{ij}|^2 \quad (12)$$

where

$$f_i(p) = \frac{d_i}{(2\pi)^3} \exp \left(\frac{1}{T} (\mu_B + \mu_S S_i - m_i - p^2/2m_i) \right) \quad (13)$$

Here μ_B and μ_S are chemical potentials, S_i is the strangeness of baryon i , d_i is a degeneracy factor, and T is the temperature. Thermal and chemical equilibrium are assumed. The above calculation yields ⁽³⁰⁾

$$N_H = \left(\frac{2\pi}{3} \right)^{1/2} \sum_{i,j} N_i N_j C_{ij}^2 (R_H/R)^3 \left[1 + R_H^2 T \frac{m_i + m_j}{3} \right]^{-3/2} \quad (14)$$

$$N_i = d_i \left(\frac{m_i T}{2\pi} \right)^{3/2} V \exp \left(\frac{1}{T} (\mu_B + \mu_S S_i - m_i) \right)$$

Here $V = 4\pi R^3/3$ is the volume of the fireball.

In Fig. 2, we display the density dependence of the ratio N_H/N_Y ($N_Y = N_\Lambda + N_{\Sigma^0}$) as a function of T , for a “stiff” equation of state (EOS). The results are very similar for a soft EOS ⁽³⁰⁾. At AGS energies, the baryon density ρ_B characteristic of central collisions is of order $2\rho_0$ ($\rho_0 \approx 1/6 \text{ fm}^{-3}$) and the temperature is expected to be larger than the value $T \approx 120 \text{ MeV}$ obtained from proton spectra at BEVALAC energies ⁽³⁴⁾. A recent analysis ⁽³⁵⁾ of transverse energy distributions in Si + Pb central collisions at the AGS is compatible with $\rho_B/\rho_0 \approx 2.7$ and an energy density 0.8 GeV/fm^3 . From Fig. 2, this implies a ratio N_H/N_Y of order 5×10^{-3} to 10^{-2} at AGS energies. Since $N_Y \approx 4$, as per Eq. (8), we predict ⁽³⁰⁾

$$N_H \approx (2 - 4) \times 10^{-2} \quad (15)$$

per central $^{28}\text{Si} + ^{197}\text{Au}$ collision. For $^{197}\text{Au} + ^{197}\text{Au}$ collisions, N_H is even larger, of order 10^{-1} . In terms of cross section σ_H for H production, we have

$$\sigma_H \approx \frac{N_H}{N_Y} \sigma_Y \approx 2 \frac{N_H}{N_Y} (A_p A_T)^{2/3} \sigma(pp \rightarrow \Lambda X) . \quad (16)$$

Using $\sigma(pp \rightarrow \Lambda X) \approx 0.7 \text{ mb}$, appropriate for a c.m. energy $\sqrt{s} \approx 4.5 \text{ GeV}$ ⁽³⁶⁾, we estimate $\sigma_H \approx 5 - 10 \text{ mb}$ for $^{28}\text{Si} + ^{197}\text{Au}$ at 15 GeV/A . This rather substantial cross section indicates that relativistic heavy ion H collisions afford a very promising way to produce the H dibaryon.

4. - Detection of the H

In (K^-, K^+) experiments, it is not necessary to detect the H directly, since the final state, for instance in $\Xi^- + d \rightarrow H + n$ or $K^- + ^3\text{He} \rightarrow H + n + K^+$ reactions, contains only two or three particles. Thus, one performs a “missing mass” experiment. In relativistic heavy ion collisions, such an approach is not feasible, since one is faced with a high multiplicity of particles (pions, protons, etc.) produced in a central collision. Thus, the neutral H particle must be detected directly. Dover, Koch and May ⁽³⁰⁾ have discussed several possible methods for detecting the H. In this chapter, we summarize these considerations.

4.1 Weak decays of the H

The only calculation of H weak decay which takes full account of the microscopic quark structure of the H wave function is due to Donoghue *et al* ⁽³²⁾. Their estimate of the H lifetime τ_H , as a function of mass, is shown in Fig. 3. This calculation has the remarkable feature that τ_H is at least a factor of 10 larger than the Λ lifetime $\tau_\Lambda = 2.7 \times 10^{-10} \text{ sec}$, even when the H lies close to the $\Lambda\Lambda$ threshold. If this is true, the possibilities for experimental detection of the H are greatly improved, since the decay products of the H will be found rather far from the production point (see later). To see the origin of this effect, we note that a Q^6 system with $J^\pi = 0^+$ occurs only for flavor $\{f\} = \{1, 27, 28\}$. Since the $B_8 B_8$ system does not couple to $\{f\} = \{28\}$, the parity-conserving decay $H \rightarrow YN$ involves a

transition $\{1\} \rightarrow \{27\}$ characterized by the weak Hamiltonian

$$H_{27}^{\Delta S=1} = \frac{G_F \sin \theta_c \cos \theta_c}{2\sqrt{2}} (C_3 O_3 + C_4 O_4) \quad (17)$$

where $\{G_F, \theta_c\}$ are the Fermi coupling constant and Cabibbo angle, and $O_{3,4}$ are chiral four-quark operators ⁽³²⁾. The term $C_3 O_3$ corresponds to an isospin change $\Delta I = 1/2$, while $C_4 O_4$ is the $\Delta I = 3/2$ part. Note that the enhanced $\{8\}$, $\Delta I = 1/2$ operator $C_1 O_1$ does not contribute, and hence the usual $\Delta I = 1/2$ rule does not hold for $H \rightarrow YN(^1S_0)$ decays. In fact, since $C_4/C_3 \approx 5$, the $\Delta I = 3/2$ transition dominates, and ΣN final states are favored over Λn . This explains the rapid rise in τ_H as we drop below the ΣN threshold (see Fig. 3). As m_H decreases below the Λn threshold, only $\Delta S = 2$ decays are possible, and τ_H becomes very long ($> 3 \times 10^5$ sec). For parity-violating decays $H \rightarrow YN(^3P_0)$, the enhanced octet part $H_8^{\Delta S=1}$ (with $\Delta I = 1/2$) does contribute. The rate for these transitions, while they should be comparable ⁽³²⁾ to $H \rightarrow YN(^1S_0)$, is not included in Fig. 3.

Experimentally, the charged decay mode $H \rightarrow \Sigma^- p$ is most favorable for detection. Neglecting phase space factors, Donoghue et al ⁽³²⁾ predict $\Gamma(\Sigma^- p) \approx \Gamma(\Sigma^0 n) \approx 3\Gamma(\Lambda n)$, so the $\Sigma^- p$ decay mode of the H will be prominent in the range $2134 < m_H < 2231$ MeV/c². The signature for the H would be a "vee" originating far from the production point. For an H produced at rest in the c.m. system for $^{28}\text{Si} + ^{197}\text{Au}$ at 15 GeV/A, the decay length of the H is about 1.5 meters for $\tau_H = 0.5 \times 10^{-8}$ sec. This is a favorable situation for detection of the $\Sigma^- p$ "vee". On the other hand, if the H mass drops below 2134 MeV/c², H decay produces one or two neutrons, and detection becomes very difficult. In this case, the techniques discussed in Chaps. 4.2 and 4.3 become more appropriate.

4.2. Dissociation of the H

Even if the H is deeply bound, one could look for the process of dissociation on a hydrogen or complex nuclear target. Shahbazian et al ⁽²⁹⁾ have in fact reported one event which is interpreted as H dissociation. The prototype reactions are

$$H + p \rightarrow \Lambda + \Lambda + p, \Xi^- + p + p. \quad (18)$$

After weak decays $\Lambda \rightarrow p\pi^-$ or $\Xi^- \rightarrow \Lambda\pi^- \rightarrow p\pi^-\pi^-$, a five-prong final state $3p + 2\pi^-$ is produced. If one is able to detect the recoil proton ⁽³⁷⁾ as well as the charged decay products, one could determine the mass of the H . The dissociation of the H into $\Sigma^+\Sigma^-$ or $\Sigma^0\Sigma^0$ is not favorable, because of the presence of two or more neutral particles in the final state.

The threshold momentum for H dissociation on a ^{12}C target is shown in Fig. 4. as a function of m_H . The average lab momentum $\langle p_H \rangle$ of the H emitted from a typical heavy ion fireball is well above threshold. To see this, we note that the total mean lab energy $\langle E_H \rangle$ is

$$\langle E_H \rangle = \gamma \left[m_H + \frac{3}{2}T + \beta (3m_H T)^{1/2} \cos \theta \right] \quad (19)$$

where

$$\begin{aligned}\gamma &= \cosh y \\ \beta &= \frac{(\gamma^2 - 1)^{1/2}}{\gamma}\end{aligned}\quad (20)$$

and y is the fireball rapidity, θ the c.m. emission angle of the H. For $^{16}\text{O} + ^{197}\text{Au}$ collisions at 15 GeV/A, we have $\gamma = 1.83$, $\beta = 0.84$. The maximum lab emission angle for the H is about 17.8° , corresponding to $\theta = 120^\circ$. For this case, we have $\langle p_H \rangle \approx \frac{3}{2}m_H$, while for $\theta = 0^\circ$ emission, the result is $\langle p_H \rangle \approx \sqrt{8}m_H$. In all cases, $\langle p_H \rangle$ is more than 3 GeV/c, so we are far above the thresholds indicated in Fig. 4.

For large $\langle p_H \rangle$, the main contribution to the cross section for the reactions of Eq. (18) comes from the dissociation of the H into baryons of momenta around $\langle p_H \rangle/2$, with low relative momentum, followed by a Λp , $\Xi^- p$ or pp final state interaction which serves to put the process on-shell. A very rough estimate yields

$$\sigma \approx \begin{cases} \frac{1}{4}\sigma_{pp} \approx 10\text{mb for } Hp \rightarrow \Xi^- pp \\ \frac{1}{8}\sigma_{\Lambda p} \approx 3\text{mb for } Hp \rightarrow \Lambda \Lambda p \end{cases}\quad (21)$$

assuming $\sigma_{pp} \approx 40$ mb and $\sigma_{\Lambda p}/\sigma_{pp} \approx 2/3$.

The presence of a large neutron flux could constitute an important background which might mask the signal due to H dissociation. The cross section for the process $np \rightarrow 5$ prongs varies from 0.2 to 0.6 mb in the 3–4 GeV/c region ⁽³⁸⁾, which is much smaller than Eq. (21). However, $N_n/N_H \approx 10^4$, so a detector would need to distinguish between the $3p2\pi^-$ final state (the signal) and events containing combinations like $2p\pi^+ 2\pi^-$, which could arise from np interactions.

4.3. Nuclear fragments containing the H

If the H is very long-lived, as would be the case if $m_H < 2055 \text{ MeV}/c^2$, it would not be practical to look for the weak decay $H \rightarrow nn$ in the debris of a heavy ion collision. An alternative strategy is to look for nuclear fragments ⁽³⁰⁾ of abnormal charge to mass ratio Q/M , consisting of an H bound to a nuclear core consisting of Z protons and N neutrons (with $A = N + Z$). Writing $m_H = (2 + x)m_n$, where $0 < x < 2$ ($m_\Lambda/m_n - 1$), and neglecting the binding energy, we have (M in units of m_n)

$$\frac{Q}{M} \approx \frac{Z}{A + 2 + x} \quad (22)$$

For the systems $H + d$ (^4H), $H + t$ (^5H) and $H + \alpha$ (^6He), we have

$$\frac{Q}{M} = \begin{cases} 0.110 - 0.125 & (^4\text{H}) \\ 0.186 - 0.2 & (^5\text{H}) \\ 0.313 - 0.33 & (^6\text{He}) \end{cases}\quad (23)$$

A few neutron-rich metastable nuclear species also have a low Q/M (ex. ${}^8\text{He}$, with $Q/M \approx 1/4$). If necessary, these could be distinguished from H-nuclei by a measurement of M .

Since the H is much more massive than the nucleon, only a rather shallow attractive potential is needed for it to bind to a nucleus. For a square well potential of depth V_0 and radius $R = r_0 A^{1/3}$, the requirement for binding the s -wave is $V_0 \geq \pi^2/8\mu R^2$, where $\mu = m_H/(1 + m_H/A m_n)$ is the reduced mass. For $r_0 \approx 1.4$ fm, the minimum value of V_0 is shown in Fig. 5 as a function of the atomic number A of the core. Even for a light core with $A = 2$, only a modest value of $V_0 = 16$ MeV suffices to bind the H; this is small compared to the well depth of 50 – 60 MeV characteristic of the nucleon-nucleus potential at low energy.

No detailed calculation of the H-nucleus well depth is possible, but one can make a crude estimate ⁽³⁰⁾ of the attractive contribution coming from second order pion exchange, i.e., $\text{HN} \xrightarrow{\pi} \text{H}'_1 \text{N} \xrightarrow{\pi} \text{HN}$, where H'_1 is the 1^+ , $I = 1$ first excited state shown in Fig. 1. This is analogous to the chains $\text{NN} \xrightarrow{\pi} \Delta \text{N} \xrightarrow{\pi} \text{NN}$ or $\Delta \text{N} \xrightarrow{\pi} \Sigma \text{N} \xrightarrow{\pi} \Delta \text{N}$, which contribute attraction to the N-nucleus or Δ -nucleus potentials, respectively. If we assume a transition potential for $\text{HN} \rightarrow \text{H}'_1 \text{N}$ of the form $V_{\text{tr}} = \alpha \exp(-\mu_\pi r)/\mu_\pi r$, then the effective second order diagonal potential is

$$V(\text{HN} \rightarrow \text{HN}) \approx -V_{\text{tr}}^2/\Delta, \quad (24)$$

where $\Delta \approx m_{\text{H}'_1} - m_H \approx 200$ MeV. Converting $V(\text{HN} \rightarrow \text{HN})$ to an equivalent zero range potential and convoluting it with the nuclear density $\rho(r)$, we obtain an H-nucleus potential

$$V_H(r) \approx -V_0 \rho(r)/\rho_0 \quad (25)$$

with $V_0 = 2\pi\rho_0\alpha^2/\mu_\pi^3\Delta$. For $\alpha \approx 10$ MeV, a strength comparable to the $\Delta \text{N} \xrightarrow{\pi} \Sigma \text{N}$ coupling, we obtain $V_0 \approx 2$ MeV. This is not sufficient by itself to bind the H, but it represents only the contribution of the central part of V_{tr} . For the $\Delta \text{N} \rightarrow \Sigma \text{N} \rightarrow \Delta \text{N}$ case, the contribution of second order tensor (S_{12}) forces is also attractive, and about an order of magnitude larger than the central part ⁽³⁹⁾. If we suppose this situation also prevails for the H, although this must be verified by a detailed calculation, a contribution to the well depth V_0 of order 20 MeV might be expected from tensor forces. In this case, even a light system like $\text{H} + d$ could be bound.

The production of H-nuclei may proceed via a second order coalescence mechanism in which $\{d, t, \alpha\}$ clusters are formed, subsequently fusing with the H to form $\{{}^4_{\text{H}}, {}^5_{\text{H}}, {}^6_{\text{H}}\}$. Bandō *et al* ⁽⁴⁰⁾ have performed such calculations for $\Lambda\Lambda$ hypernuclear formation. there will be a substantial penalty factor for such a higher order fusion process; in ref. ⁽³⁰⁾, a rate of order 10^{-4} is estimated for ${}^4_{\text{H}}$ formation, relative to that for the H itself. Thus such objects will be rare, but their signature of an anomalously low Q/M ratio is quite distinctive.

Another rather speculative possibility is that of an $\text{H}^2 = (\text{HH})_{L=0}$ bound state [$J^\pi = 0^+$, $I = 0$], a $S = -4$ analog to the α particle. Only a small amount of attraction is

required to bind such an object. Some of this would be provided by second order pion exchange ($HH \xrightarrow{\pi} H'_1 H'_1 \xrightarrow{\pi} HH$, analogous to $NN \rightarrow \Delta\Delta \rightarrow NN$). In heavy ion collisions, such an object could be formed by a second order fusion process, with a rate estimated ⁽³⁰⁾ at 10^{-6} to 10^{-7} per central collision at AGS energies. The weak decays of the HH system could provide spectacular experimental signatures, for instance $H^2 \rightarrow 4\Lambda \rightarrow 4p + 4\pi^-$ (an eight-prong event) or $H^2 \rightarrow \Sigma^- p H \rightarrow \Sigma^- + 3p + 2\pi^-$ (a six-prong event).

5. - Conclusions

The search for dibaryon resonances and bound states is an exciting quest. Only in the strange sector could such objects be long-lived or even stable against strong decay. This makes the development of intense K^- beams in the 1.5–2 GeV/c region a particularly timely endeavour. The high momentum K^- beam line under construction at the Brookhaven AGS has as its highest priority the discovery of the H dibaryon (via the (K^-, K^+) reaction). This facility will be useful for other spectroscopic studies of $S = -2$ systems (Ξ or $\Lambda\Lambda$ hypernuclei, for instance). The existence of intense kaon beams at a future hadron facility would lead to even more dramatic advances in strange particle nuclear physics.

In this paper, an alternative possibility for producing strange dibaryons via relativistic heavy ion collisions was developed in detail. In such collisions at AGS energies (15 GeV/A), H 's are produced rather copiously at the level of 10^{-2} per central collision (~ 10 mb cross section). The problem is to detect the neutral H in a high multiplicity environment. Several suggestions were advanced, including the detection of weak decay products of the H , diffraction dissociation of the H , and H -nucleus bound states of anomalous charge/mass ratio. In view of our lack of knowledge of the H mass, several complementary experiments are required.

The submitted manuscript has been authored under contract DE-AC02-76CH00016 with the U. S. Department of Energy. Accordingly, the U. S. Government retains a non-exclusive, royalty-free license to publish or reproduce the published form of this contribution or allow others to do so, for U. S. Government purposes. I would like to acknowledge numerous fruitful discussions with Peter Koch and Morgan May. Most of the material presented here on H dibaryon production in heavy ion collisions is taken from our joint work (Ref. 30).

DISCLAIMER

This report was prepared as an account of work sponsored by an agency of the United States Government. Neither the United States Government nor any agency thereof, nor any of their employees, makes any warranty, express or implied, or assumes any legal liability or responsibility for the accuracy, completeness, or usefulness of any information, apparatus, product, or process disclosed, or represents that its use would not infringe privately owned rights. Reference herein to any specific commercial product, process, or service by trade name, trademark, manufacturer, or otherwise does not necessarily constitute or imply its endorsement, recommendation, or favoring by the United States Government or any agency thereof. The views and opinions of authors expressed herein do not necessarily state or reflect those of the United States Government or any agency thereof.

Table 1. Wave function components and energetically allowed BB decays for $S = -2$ Q^6 dibaryons.

State	BB Components ($L = 0$)	Allowed Decays
H	$\Lambda\Lambda, \Xi N, \Sigma\Sigma$	none
H'_0	$\Xi N, \Xi^* N, \Sigma^* \Sigma$	$\Xi N (^3S_1)$
H'_1	$\Lambda\Sigma, \Xi N, \Sigma\Sigma,$ $\Lambda\Sigma^*, \Xi^* N, \Sigma\Sigma^*, \Xi\Delta$	$\Xi N, \Lambda\Sigma (^3S_1)$
H''_0	$\Xi^* N, \Sigma^* \Sigma$	$\Lambda\Lambda, \Sigma\Sigma (^1D_2), \Xi N (^1D_2, ^3D_2)$
H''_1	$\Lambda\Sigma^*, \Xi^* N, \Sigma\Sigma^*, \Xi\Delta$	$\Sigma\Sigma (^1D_2), \Lambda\Sigma, \Xi N (^1D_2, ^3D_2)$

Table 2. Occurrence of states* $J^\pi, I, \{f\}$ in the Six Quark (Q^6) or Skyrme (12) Spectrum for $S = -2$.

$\{f\}$	J^π	I	Q^6	Skyrme	Decay to $B_8 B_8$?
1	0^+	0	yes	no	yes
8	0^+	0,1	no	no	yes
10	0^+	1	no	no	no
10^*	0^+	1	no	yes	no
27	0^+	0,1	yes	yes	yes
8	1^+	0,1	yes	no	yes
10	1^+	1	yes	no	yes
10^*	1^+	1	yes	yes	yes
27	1^+	0,1	no	yes	no

Note that the combinations $\{f\}, I = \{1\}, 1; \{10\}, 0; \{10^\}, 0$ are excluded for hypercharge $Y = 0$ by general symmetry arguments.

References

1. R. L. JAFFE: *Phys. Rev. Lett.*, 38, 195 (1977) and 38, 1617E (1977).
2. P. J. G. MULDER, A. T. M. AERTS and J. J. DE SWART: *Phys. Rev.*, D21, 2653 (1980).
3. P. J. G. MULDER and A. W. THOMAS: *J. Phys.*, G9, 1159 (1983).
4. A. T. M. AERTS and J. RAFELSKI: *Phys. Lett.*, 148B, 337 (1984).
5. B. O. KERBIKOV: *Yad. Fiz.*, 39, 816 (1984) [*Sov. J. Nucl. Phys.*, 39, 516 (1984)].
6. J. L. ROSNER: *Phys. Rev.*, D33, 2043 (1986).
7. A. P. BALACHANDRAN, A. BARDUCCI, F. LIZZI, V. G. J. RODGERS and A. STERN: *Phys. Rev. Lett.*, 52, 887 (1984).
8. A. P. BALACHANDRAN, F. LIZZI, V. G. J. RODGERS and A. STERN: *Nucl. Phys.*, B256, 525 (1985).
9. R. L. JAFFE and C. L. KORPA: *Nucl. Phys.*, B258, 468 (1985).
10. S. A. YOST and C. R. NAPPI: *Phys. Rev.*, D32, 816 (1985).
11. C. G. CALLAN and I. KLEBANOV: *Nucl. Phys.*, B262, 365 (1985).
12. J. KUNZ and P. J. MULDER: NIKHEF preprint, Amsterdam (July 1988).
13. C. B. DOVER: *Nucl. Phys.*, A450, 95c (1986).
14. P. B. MACKENZIE and H. B. THACKER: *Phys. Rev. Lett.*, 55, 2359 (1985).
15. Y. IWASAKI, T. YOSHIE and Y. TSUBOI: *Phys. Rev. Lett.*, 60, 1371 (1988).
16. M. OKA, K. SHIMIZU and K. YAZAKI: *Phys. Lett.*, 130B, 365 (1983).
17. U. STRAUB, ZONG-YE ZHANG, K. BRÄUER, A. FAESSLER and S. B. KHADKIKAR: *Phys. Lett.*, 200B, 241 (1988).
18. M. BOZOIAN, J. C. H. VAN DOREMALEN and H. J. WEBER: *Phys. Lett.*, 122B, 138 (1983).
19. T. H. R. SKYRME: *Proc. Roy. Soc. London*, Ser. A260, 127 (1961) and *Nucl. Phys.*, 31, 556 (1962).
20. E. WITTEN: *Nucl. Phys.*, B223, 422 and 433 (1983).
21. P. J. MULDER, A. T. M. AERTS and J. J. DE SWART: *Phys. Rev. Lett.*, 40, 1543 (1978).
22. G. FRANKLIN et al: Brookhaven AGS experiments 813 and 835.
23. K. IMAI et al: KEK (Japan) proposal.
24. A. T. M. AERTS and C. B. DOVER: *Phys. Rev.*, D29, 433 (1984).
25. A. T. M. AERTS and C. B. DOVER: *Phys. Rev. Lett.*, 49, 1752 (1982) and *Phys. Rev.*, D28, 450 (1983).

26. P. BARNES: these proceedings.
27. K. KILIAN: in *Physics at LEAR with Low Energy Antiprotons*, LEAR Workshop IV, Villars-sur-Ollon, September 1987, Nuclear Science Research Conf. Series, Eds. C. Amsler et al, Harwood Academic Publishers, Chur, Vol. 14, 529 (1988).
28. C. GUARALDO: these proceedings.
29. B. A. SHAHBAZIAN, A. O. KECHECHYAN, A. M. TARASOV and A. S. MARTYNOV: *Z. Phys.*, C39, 151 (1988).
30. C. B. DOVER, P. KOCH and M. MAY: Brookhaven preprint (1988).
31. S. NAGAMIYA: private communication.
32. J. F. DONOGHUE, E. GOLOWICH and B. R. HOLSTEIN: *Phys. Rev.*, D34, 3434 (1986).
33. P. GONZALEZ and V. VENTO: *Few-Body Systems*, 2, 145 (1987).
34. H. STÖCKER and W. GREINER: *Phys. Rep.*, 137, 277 (1986).
35. B. BRAUN-MÜNZINGER et al: Brookhaven preprint BNL 40838 (1988) and *Z. Phys. C* (1988).
36. V. FLAMINIO et al: CERN-HERA report 84-01 (1984).
37. S. NAGAMIYA: private communication.
38. A. ABDIVALIEV et al: *Nucl. Phys.*, B99, 445 (1975).
39. C. B. DOVER and A. GAL: *Prog. in Particle and Nuclear Phys.*, 12, 171 (1984).
40. H. BANDŌ, M. SANO, M. WAKAI and J. ZOFKA: Institute of Nuclear Studies (Tokyo) preprint INS-Rep-682 (1988).

Fig. 1. The spectrum of $S = -2$ six-quark states, from Mulders et al ⁽²⁾. Various strong decay thresholds are indicated by dashed lines.

Fig. 2. The ratio N_H/N_Y as a function of baryon density ρ_B for various values of temperature T , from ref. ⁽³⁰⁾.

Fig. 3. The lifetime τ_H of the H dibaryon, as a function of its mass m_H , from Donoghue *et al* (32).

Fig. 4. Threshold momentum for H dissociation into $\Lambda\Lambda$ on a ^{12}C target, as a function of the H mass. The solid curves correspond to various configurations of the recoiling $A = 12$ nucleus from the ground state of ^{12}C to the fully dissociated $6n + 6p$ system.

Fig. 5. Well depth V_0 required to bind the H to a nuclear core, as a function of A .